

INSTANTONS AND NUCLEON MAGNETISM

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We construct improved QCD sum rules for the nucleon magnetic moments by implementing direct-instanton contributions to the operator product expansion of the nucleon correlator in a magnetic background field. The instanton contributions are found to affect only those sum rules which had previously been considered unstable. The new sum rules show a high degree of stability and reproduce the experimental values of the nucleon magnetic moments for values of the magnetic quark condensate susceptibility which are consistent with other estimates. (Invited talk given at “Hadron Physics 2000”, Caraguatatuba, São Paulo, Brazil (April 10-15, 2000).

Over the last years, an increasingly multifaceted picture of the role of QCD instantons¹ in hadron physics has begun to emerge from hadron models with more or less instanton-induced structure², instanton-liquid vacuum models³, a generalized operator product expansion including instanton contributions (IOPE)⁴ and, most recently, lattice calculations⁵.

One of the major benefits of sharpening this picture by identifying specific hadronic instanton effects will be to end the puzzling elusiveness of explicit glue in hadrons at low energies. Indeed, in most successful models of the classical mesons and baryons, gluons play at best a marginal role. The situation is further obfuscated by the fact that gluons do not directly couple to electroweak probes. Since much of what we know about the classical hadrons derives from their response to such probes, it is particularly desirable to pin down traces of nonperturbative glue in electromagnetic hadron observables.

In this talk, I will report on an investigation⁶ which takes a step towards uncovering the hidden role of glue in the response to electromagnetic fields by studying the impact of instantons on the nucleon magnetic moments in the recently developed IOPE sum-rule approach^{4,7}. While in general the range of applicability of the IOPE (as a short-distance expansion) is more limited than that of instanton vacuum models or lattice calculations, it profits from the transparency of an analytical (yet largely model-independent) method and takes, in contrast to instanton models, all long-wavelength vacuum fields and also perturbative fluctuations into account.

In order to study the nucleon’s magnetic moments in the framework of the IOPE, we start from a correlation function which contains the response of the nucleon to the electromagnetic current at short, spacelike distances. (The

information on the magnetic moments is subsequently extracted by dispersion techniques^{4,8}.) At first sight, the natural candidate for such a correlation function would be

$$\Pi_\mu(p, q) = - \int d^4x e^{ipx} \int d^4y e^{iqy} \langle 0 | T \eta(x) J_\mu(y) \bar{\eta}(0) | 0 \rangle \quad (1)$$

where the nucleon interpolating fields $\eta(x)$ (with either proton or neutron quantum numbers) are composite operators of massless up and down quark fields⁹,

$$\eta_p(x) = \epsilon^{abc} [u^{aT}(x) C \gamma_\alpha u^b(x)] \gamma_5 \gamma^\alpha d^c(x), \quad \eta_n = \eta_p(u \leftrightarrow d), \quad (2)$$

and J_μ is the electromagnetic current. However, obtaining the nucleon magnetic moments requires taking the $q \rightarrow 0$ limit in which the matrix element of (1) is probed mainly in the large-distance region $|y|, |x - y| \gg \Lambda^{-1}$ where the operator product expansion breaks down^a.

Fortunately, there exists an alternative approach, due to Ioffe and Smilga¹¹ and Balitsky and Yung¹², which can be directly applied to the calculation of the magnetic moments. The underlying idea is to treat the long-wavelength photons at the same level as the soft QCD vacuum fields (which make up the condensates), i.e. to account for them in the operators of the (I)OPE (see below). This amounts to considering the electromagnetic field as a constant, classical background field (which, incidentally, is reminiscent of the experimental setup used in actual measurements of the moments). In other words, the electromagnetic probe does not appear anymore as an external field but rather as part of the “magnetized” vacuum state.

We are thus led to consider the nucleon correlation function

$$\Pi(p) = i \int d^4x e^{ipx} \langle 0 | T \eta(x) \bar{\eta}(0) | 0 \rangle_F \quad (3)$$

$$= \Pi_0(p) + \sqrt{4\pi\alpha} \Pi_{\mu\nu}(p) F^{\mu\nu} + O(F^2) \quad (4)$$

in the background of a constant electromagnetic field $F_{\mu\nu}$. We have not written down higher orders of F since the information on the magnetic moments is part of the linear response

$$\Pi_{\mu\nu}(p) = (\not{p} \sigma_{\mu\nu} + \sigma_{\mu\nu} \not{p}) \Pi_1(p^2) + i(\gamma_\mu p_\nu - \gamma_\nu p_\mu) \not{p} \Pi_2(p^2) + \sigma_{\mu\nu} \Pi_3(p^2) \quad (5)$$

to the (arbitrarily weak) external field. Above we have given the general decomposition of $\Pi_{\mu\nu}(p)$ in terms of three independent Lorentz and spinor structures which are associated with one chirally-even (Π_1) and two chirally-odd

^a The OPE of this matrix element can be used, however, to obtain the nucleon form factors at intermediate momentum transfers $Q \sim 1 \text{ GeV}$ ¹⁰.

invariant amplitudes (Π_2 and Π_3). The relation to the magnetic moments can be made explicit by writing

$$\Pi(p) = i \int d^4x e^{ipx} \langle 0 | T \eta(x) \bar{\eta}(0) e^{-i \int d^4y A^\mu J_\mu} | 0 \rangle \quad (6)$$

where $J_\mu = \bar{q} \gamma_\mu Q q$ is the electromagnetic current with quark charge matrix Q and A_μ is the external vector potential. In fixed-point gauge, and specializing to a constant magnetic background field

$$B_i = -\frac{1}{2} \varepsilon_{ijk} F_{jk}, \quad (E_i = F_{i0} = 0) \quad (7)$$

with the potential

$$A_\mu(y) = -\frac{1}{2} F_{\mu\nu} y^\nu, \quad (8)$$

the above exponential can be rewritten as

$$\int d^4y A^\mu J_\mu = \frac{1}{2} F_{ij} \int d^4y y_j J_i = \int dt \vec{\mu} \cdot \vec{B} \quad (9)$$

where $\vec{\mu}$ is the magnetic-moment operator

$$\vec{\mu} = \frac{1}{2} \int d^3x (\vec{x} \times \vec{J}). \quad (10)$$

Expanding (6) to first order in F and inserting nucleon intermediate states, we thus obtain

$$\Pi_{\mu\nu}(p) F^{\mu\nu} = -\varepsilon_{ijk} \Pi_{ij}(p) B_k \sim \langle N(p) | \vec{\mu} | N(p) \rangle \cdot \vec{B} \quad (11)$$

in terms of the nucleon magnetic moments, as anticipated.

In order to write down sum rules for $\langle N(p) | \vec{\mu} | N(p) \rangle$, we need a QCD description of the correlator (3) at momenta $s = -p^2 \simeq 1 \text{ GeV}^2$, i.e. at distances $x \sim 0.2 \text{ fm}$. Such a description is provided by the nonperturbative IOPE which factorizes (3) into contributions from soft ($k_i < \nu$) and hard ($k_i \geq \nu$) field modes (with momenta k_i), where $\nu \sim 0.5 \text{ GeV}$ is the operator renormalization scale. The IOPE is generated by splitting each diagram contributing to (5) in all possible ways into a hard and a soft subgraph. The hard subgraphs, with the integration range of each internal momentum restricted^b to be larger than ν , give rise to the Wilson coefficients and receive, beyond the standard perturbative contributions¹¹, direct instanton contributions which we will evaluate

^b In practice, this restriction is often unnecessary.

below. The soft subgraphs yield hadron-channel independent condensates, i.e. vacuum expectation values of local, composite QCD operators renormalized at ν .

The new feature brought in by the magnetic background field is that the “magnetized” vacuum state $|0\rangle_F$ singles out a preferred direction (that of the field strength \vec{B}) and therefore ceases to be a Lorentz scalar. Hence, several Lorentz-covariant operators of the OPE (considered here up to dimension eight) acquire finite expectation values, the F -induced condensates

$$\langle 0|\bar{q}\sigma_{\mu\nu}q|0\rangle_F = \sqrt{4\pi\alpha}\chi F_{\mu\nu}\langle 0|\bar{q}q|0\rangle, \quad (12)$$

$$g\langle 0|\bar{q}G_{\mu\nu}q|0\rangle_F = \sqrt{4\pi\alpha}\kappa F_{\mu\nu}\langle 0|\bar{q}q|0\rangle, \quad (13)$$

$$g\langle 0|\bar{q}\gamma_5\tilde{G}_{\mu\nu}q|0\rangle_F = \frac{i}{2}\sqrt{4\pi\alpha}\xi F_{\mu\nu}\langle 0|\bar{q}q|0\rangle. \quad (14)$$

($G_{\mu\nu} = \frac{1}{2}\lambda_a G_{\mu\nu}^a$, $\tilde{G}_{\mu\nu} = \frac{1}{2}\varepsilon_{\mu\nu\rho\sigma}G_{\rho\sigma}$ with $\varepsilon_{0123} = -1$.) The parameters χ , κ , and ξ quantify the vacuum response to weak electromagnetic fields and thus play the role of generalized susceptibilities. The magnetic susceptibility of the quark condensate, χ , for example, originates from the induced spin alignment of quark-antiquark pairs in the vacuum. The corresponding vacuum expectation value is the lowest-dimensional induced condensate and will therefore play a leading role in the magnetic moment sum rules.

Altogether, the IOPE of the invariant amplitudes $\Pi_i(p)$ in Eq. (5) can thus be written as

$$\Pi_i(p^2) = \sum_n C_{i,n}(p^2; \nu) \langle 0|\mathcal{O}_n[\nu]|0\rangle_F \quad (15)$$

where the \mathcal{O}_n are local, composite QCD operators with the appropriate quantum numbers (the multi-index n can contain Lorentz indices) and dimension $d \leq 8^c$. The C_n are the Wilson coefficients. It might be worth reiterating that the instanton contributions to (15) factorize into those from soft modes (contained in the condensates $\langle 0|\mathcal{O}_n|0\rangle$) and those due to hard modes which add nonperturbative structure to the C_n . The latter can be of substantial size since $\bar{\rho} < \nu^{-1} \simeq 0.4\text{ fm}$ and have to be calculated explicitly. (Nevertheless, they were neglected in the usual, perturbative treatment.)

The calculation of the IOPE coefficients is described in Ref. ^{4,6} to which we refer for details. Similar to the phenomenologically known values of the lowest-dimensional condensates, the bulk properties of the instanton size distribution are generated by long-distance vacuum dynamics and thus taken as input. We

^c Contributions from higher-dimensional operators are strongly suppressed, see below.

will use the standard “instanton liquid” values³ $\bar{\rho} \simeq \frac{1}{3} \text{ fm}$ for the average instanton size and $\bar{R} \simeq 1 \text{ fm}$ for the average separation between neighboring (anti)instantons. Since the instanton size distribution $n(\rho)$ is rather sharply peaked around $\bar{\rho}$, we can thus approximately set $n(\rho) = \bar{n}\delta(\rho - \bar{\rho})$ with $\bar{n} \simeq 0.5 \text{ fm}^{-4}$. Moreover, $\bar{\rho}^{-1} \gg \Lambda_{QCD}$ implies that the hard instanton contributions can be calculated semiclassically. Since, finally, multi-instanton correlations are negligible at the relevant distances $x \sim 0.2 \text{ fm} \ll \bar{R}$, this amounts to evaluating the hard subgraphs in the background of an (anti-) instanton field and subsequently averaging over the instanton-parameter distribution. The main contributions typically arise from the quark zero modes in the instanton field (which acquire an effective mass¹³ $\bar{m}(\rho) = -\frac{2}{3}\pi^2\rho^2\langle\bar{q}q\rangle$ due to the interactions with ambient vacuum fields) and can be treated exactly, while the continuum modes are approximated by plane waves.

More specifically, the leading instanton contributions to the correlator (3) can be traced to the graph in which two of the quarks propagate in zero modes while the third interacts with the background field through the magnetized quark condensate^d. This contribution is purely nonperturbative and difficult to account for, e.g., in quark-based hadron models¹⁵. After averaging over the instanton size distribution, it reads

$$\langle 0|T\eta_p(x)\bar{\eta}_p(0)|0\rangle_{F,inst} = \Pi_0^{inst}(x) - \frac{2^3 e_u}{3\pi^4} \frac{\bar{\rho}^4}{\bar{m}^2} \times \langle \bar{q}\sigma_{\mu\nu}q\rangle_F \sigma_{\mu\nu} \int d^4x_0 \frac{1}{(r^2 + \bar{\rho}^2)^3 (x_0^2 + \bar{\rho}^2)^3} \quad (16)$$

for the proton. Π_0^{inst} contains the instanton contributions for $F = 0$, calculated in⁷, and $r = x - x_0$, where x_0 specifies the center of the instanton. The corresponding neutron correlator is obtained by replacing e_u with e_d . Above, we have used the self-consistency condition¹⁶

$$\langle \bar{q}q \rangle = -2 \int d\rho \frac{n(\rho)}{\bar{m}(\rho)} = -2 \frac{\bar{n}}{\bar{m}(\bar{\rho})} \quad (17)$$

to eliminate the \bar{n} dependence from (16). For later use in the sum rules, we now calculate the Fourier and Borel transform^e of (16), from which we obtain

$$\hat{\Pi}_3(M^2) = \frac{e_u}{128\pi^4} a\chi\bar{\rho}^2 M^6 I(z^2) \quad (18)$$

^dBackground-field induced transitions between zero and continuum modes (which can be dominant in other correlators¹⁴) are absent here.

^e For the Borel transform we follow the convention of Ref.⁹.

where M denotes the Borel mass parameter and $z = M\bar{\rho}$, $a = -(2\pi)^2 \langle \bar{q}q \rangle$. The zero-mode loop gives rise to the integral

$$I(z^2) = \int_0^1 \frac{d\alpha}{\alpha^2(1-\alpha)^2} e^{-\frac{z^2}{4\alpha(1-\alpha)}} = 4e^{-\frac{z^2}{2}} \left[K_0\left(\frac{z^2}{2}\right) + K_1\left(\frac{z^2}{2}\right) \right] \quad (19)$$

and exhibits the unique exponential Borel-mass dependence which is characteristic for instanton contributions⁷.

A first important lesson of Eq. (18) is that, to leading order, direct instantons contribute almost exclusively^f to the chirally-odd amplitude Π_3 . This is crucial because in the analysis of Ref. ¹¹, where this contribution was not accounted for, exactly the Π_3 sum-rule failed to show a fiducial stability region¹⁷. The previous neglect of the instanton contributions thus offers a potential explanation for this instability. An additional hint in this direction comes from two other chirally-odd nucleon sum rules^{7,15} where stabilization due to instantons was indeed found to take place⁹.

We are now ready to proceed to the quantitative analysis of the IOPE sum rule for Π_3 . The latter results from equating the Borel transform of the standard OPE of Ref. ¹¹ and the instanton contributions (18) to the Borel-transformed double dispersion relation for the correlator (3), with a spectral function parametrized in terms of the nucleon pole contribution (cf. Eq. (11)) and a continuum based on local duality. Including the infrared-divergent term encountered in Ref. ¹¹, duly truncated at the OPE renormalization point, the IOPE sum rule (for the proton) reads

$$\begin{aligned} & aM^2 \left\{ \left[e_u - \frac{1}{6}e_d(1 + 4\kappa + 2\xi) \right] E_1(M) \right. \\ & + \frac{1}{6}e_u \frac{m_0^2}{M^2} \left[\ln \frac{M^2}{\nu^2} - \gamma_{EM} \right] L^{-\frac{4}{9}} + \frac{1}{6}e_d M^2 \chi E_2(M) L^{-\frac{16}{27}} \\ & \left. - \frac{1}{8}e_u \chi \rho_c^2 M^4 I(z^2) L^{-\frac{16}{27}} \right\} \\ & = \frac{1}{4} \tilde{\lambda}_N^2 m e^{-\frac{m^2}{M^2}} \left[\frac{\mu_p}{M^2} - \frac{\mu_p^a}{2m^2} + A_p \right], \end{aligned} \quad (20)$$

where m is the nucleon mass, μ_p (μ_p^a) the proton's (anomalous) magnetic moment, W the continuum threshold, and λ_N the coupling of the current (2) to

^f The small direct-instanton contributions to Π_2 have no appreciable impact on the sum rules and are neglected.

^g A previously contemplated connection between the instability of the original Π_3 sum rule and the perturbative infrared singularities of the “pragmatic” OPE seems, in view of newer results¹⁸, unlikely⁶.

the nucleon state, $\langle 0|\eta|N\rangle = \lambda_N u$. For the mixed quark condensate we use the standard parametrization $\langle \bar{q}\sigma_{\mu\nu}G^{\mu\nu}q\rangle = -m_0^2\langle \bar{q}q\rangle$ with $m_0^2 = 0.8\text{ GeV}^2$, and $\gamma_{EM} \simeq 0.577$ is the Euler-Mascheroni constant. The additional parameters $A_{p,n}$ determine the strength of electromagnetically induced transitions between the nucleon and its excited states. The sum rule for the neutron is obtained from (20) by interchanging e_u and e_d and by replacing $\mu_p, \mu_p^a \rightarrow \mu_n$ and $A_p \rightarrow A_n$. We have also defined $\tilde{\lambda}_N^2 = 32\pi^4\lambda_N^2$, $L = \ln(M/\Lambda)/\ln(\mu/\Lambda)$ ($\Lambda = 0.1\text{ GeV}$), and transferred, using the standard expressions

$$E_n(M) = 1 - e^{-\frac{W^2}{M^2}} \left[1 + \sum_{j=1}^n \frac{1}{j} \left(\frac{W^2}{M^2} \right)^j \right], \quad (21)$$

the continuum contributions to the IOPE-side of the sum rules. The appropriate form of these contributions has recently been clarified in Ref. ¹⁹.

The quantitative analysis of the IOPE sum rule (20) is based on the minimization of the relative deviation between its two sides. Since the fiducial domain, i.e. the Borel-mass region in which the IOPE truncation is justified while the nucleon pole still dominates over the continuum, is (as usual) not large enough to determine all unknown parameters from a combined fit, we follow the procedure of Ref. ¹¹ and obtain the coupling $\tilde{\lambda}^2 = 2.93\text{ GeV}^6$ and the continuum threshold $W = 1.66\text{ GeV}$ by fitting the instanton-improved nucleon mass sum rule of Ref. ⁷ to the experimental nucleon mass. The values of the two susceptibilities $\kappa = -0.34 \pm 0.1$ and $\xi = -0.74 \pm 0.2$ were estimated in independent work by Kogan and Wyler²⁰. This enables us to fit both sides of the sum rules (20) by varying χ and A_p (or A_n , respectively) while keeping the magnetic moments fixed at their experimental values.

The fits are performed in the fiducial Borel mass domain $0.8\text{ GeV} \leq M \leq 1.15\text{ GeV}$, where the highest-dimensional operators contribute at most 10 % to the IOPE and the continuum contribution does not exceed 50 %. This fiducial domain is larger, incidentally, than that of the sum rules based on Π_1 and Π_2 .

Figure 1 shows, both for the proton and the neutron sum rules, the direct-instanton contributions, the remaining OPE including the continuum contributions, their sum (which makes up the left-hand side of Eq. (20)), and the right-hand sides. The fit quality of both the proton and neutron sum rules is quite impressive. Figure 1 shows, furthermore, that the direct-instanton contributions can reach the magnitude of the remaining terms in the OPE, which explains why their previous neglect had a detrimental impact on the sum-rule stability. In Fig. 2 the optimized sum rules are solved for μ_N and plotted as a function of the Borel mass. The resulting functions $\mu_{p,n}(M)$ therefore specify the value of the magnetic moment which is required to make both sides of the

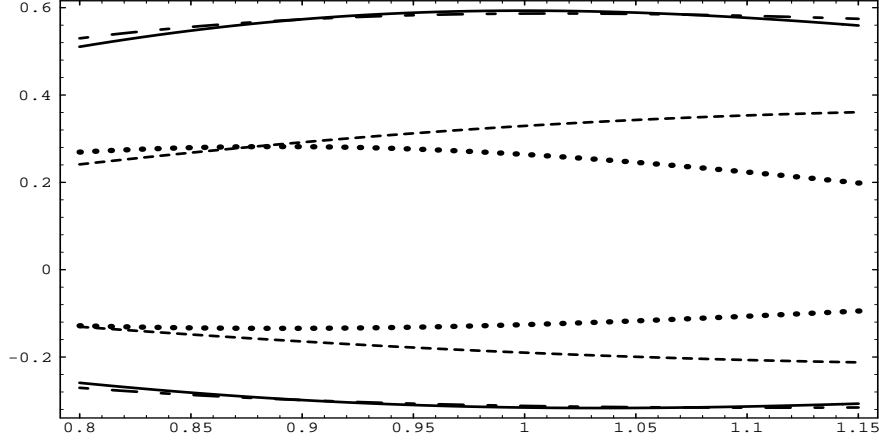


Figure 1: The OPE (dashed line) and direct instanton (dotted line) contributions to the new $\sigma_{\mu\nu}$ sum rules for the proton (positive range) and neutron. Their sum (dot-dashed line) is compared to the RHS (solid line).

sum rule (20) match exactly at a given value of M . As a consequence of the high fit quality, $\mu(M)$ is practically M -independent. The numerical results are $\chi \simeq -4.96 \text{ GeV}^{-2}$, $A_p \simeq 0.28 \text{ GeV}^2$ for the proton and $\chi \simeq -4.73 \text{ GeV}^{-2}$, $A_n \simeq -0.27 \text{ GeV}^{-2}$ for the neutron sum rule. The values of the quark condensate susceptibility χ (at $\nu = 0.5 \text{ GeV}$) lie within the range obtained from other estimates^{11,18,21} and differ somewhat from the value $\chi \simeq -5.7 \text{ GeV}^{-2}$ found in the two- and three-pole models of Ref.²¹.

In conclusion, we have recovered a third reliable sum rule for the nucleon magnetic moments. In contrast to the other two, it receives previously neglected direct-instanton contributions which arise from the interplay with long-wavelength vacuum fields. Our new sum rule is built on the chirally-odd amplitude Π_3 of the nucleon correlator in an electromagnetic background field and found to be at least as stable as the other two (although it had previously been regarded as flawed). The new sum rule adds to the predictive power of the background-field sum rules and strengthens their mutual consistency.

Moreover, our results reinforce a systematic pattern which emerged from the study of direct-instanton effects in the pion¹⁴, nucleon^{7,15} and glueball²² channels: those sum rules which worked satisfactorily without instanton corrections receive little or no direct instanton contributions while previously

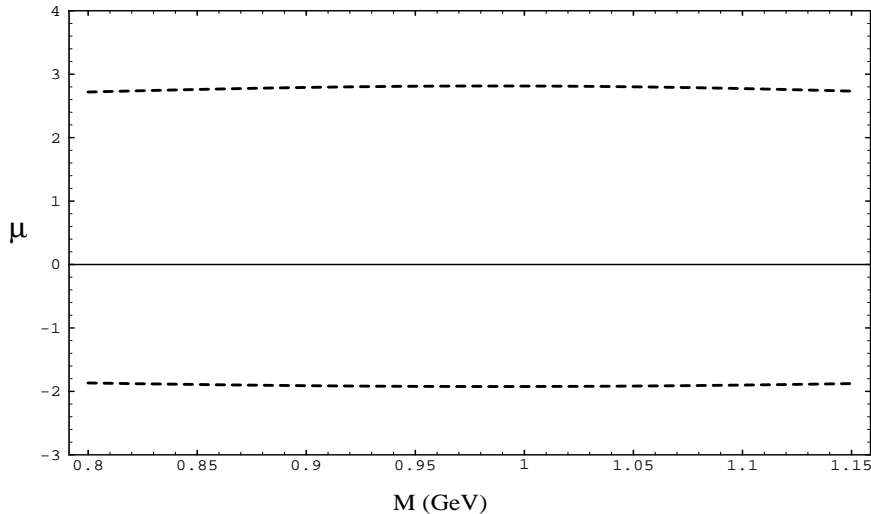


Figure 2: The Borel mass dependence of the magnetic moments of the proton (upper) and neutron calculated from the optimal fit of LHS and RHS.

less reliable or completely unstable sum rules are stabilized by large instanton contributions. This pattern points not only towards the importance of direct instantons in particular sum rules, but also supports the adequacy of their semiclassical implementation into the OPE. Our results show that these conclusions continue to hold in the presence of a “magnetized” vacuum.

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